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# The Collisional Drift Mode in a Partially Ionized Plasma

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## Abstract

The structure of the drift instability is examined in several density regimes. Let $\lambda_e$ be the total electron mean free path, $k_z$ the wave-vector component along the magnetic field, and $\nu_\perp/\nu_\parallel$ the ratio of perpendicular ion diffusion to parallel electron streaming rates. At low densities ($k_z\lambda_e > 1$) the drift mode is isothermal and should be treated kineticly. In the finite heat conduction regime $\sqrt{m/M} < k_z\lambda_e < 1$ the drift instability threshold is reduced at low densities ($\nu_\perp/\nu_\parallel < 0.1$) and increased at high densities ($\nu_\perp/\nu_\parallel > 0.1$) as compared to the isothermal threshold. Finally, in the energy transfer limit ($k_z\lambda_e < \sqrt{m/M}$) the drift instability behaves adiabatically in a fully ionized plasma and isothermally in a partially ionized plasma for an ion-neutral to Coulomb collision frequency ratio $\nu_{in}/\nu_{ii} > 2\sqrt{m/M}$ at $T_i = T_e = T_n$.

## Introduction

The most obvious source of free energy for the instability of the nighttime equatorial F region called spread F is the sharp plasma density gradient on the underside of the postsunset F layer. While a variety of density gradient dependent modes have been considered (cf. Farley et al. 1970) to explain the geomagnetic field alligned density irregularities with perpendicular wavelengths from a few m to km, Hudson and Kennel (1974a) are the first to consider collisional drift waves driven solely by the density gradient. The collisional drift mode appeared in the electrostatic limit of the electromagnetic dispersion relation for low frequency density-gradient driven modes in a partially ionized collisional plasma derived in the companion paper (Hudson and Kennel 1974b). The observed altitude threshold for spread F onset (Farley et al. 1970) suggested the extension of previous work on the collisional drift mode (Tsai et al. 1970) to partially ionized plasmas where neutral collisions affect the instability threshold and growth rate.

Recent work on the collisional drift mode has shown that its growth rate is modified by finite electron heat conduction along the magnetic field at parallel wavelengths exceeding the electron momentum transfer mean free path $\lambda_e$. However, there is an apparent paradox in the literature; Ellis and Motley (1974) find that the drift mode growth rate is enhanced by finite heat conduction, while Tsai et al. (1970) conclude that it is reduced, especially at high densities. We resolve this conflict by obtaining an explicit criterion which separates the finite heat conduction regime $\sqrt{m/M} < k_z\lambda_e < 1$ into low and high density regions where the drift mode growth

rate is enhanced or reduced from its isothermal value. A single parameter which is the ratio of the perpendicular ion diffusion to parallel electron streaming rates, $\nu_\perp/\nu_\parallel = 0.1$, separates the low and high density regions. In application to equatorial F region irregularities $\nu_\perp/\nu_\parallel < 0.1$ in general (Hudson and Kennel 1974a) and the drift mode growth rate is enhanced by finite heat conduction.

Since equatorial F region geometry permits very long parallel wavelengths, we have extended the analysis of the collisional drift mode to parallel wavelengths exceeding the energy transfer mean free path, $k_z\lambda_e < \sqrt{m/M}$, where we find that temperature fluctuations and finite heat conduction effects are surpressed. In this long wavelength regime which we call the energy transfer limit, a fully ionized plasma behaves adiabatically, because there is a balance of energy exchange between electrons and ions. A partially ionized plasma with $\nu_{in}/\nu_{ii} > 2\sqrt{m/M}$ at $T_e = T_i = T_n$ behaves isothermally (Hudson 1974), because energy in electron temperature fluctuations is lost to the ions and then lost to the neutral sink faster than it can be transferred back to the electrons.

## Electrostatic Dispersion Relation

Finite electron heat conduction, energy transfer between species and neutral collisions have all been included in our derivation of the electrostatic dispersion relation for the collisional drift mode from two fluid theory in the companion paper (Hudson and Kennel 1974b)

$$b\left\{\left(\omega + R_T\omega_D\right)(\omega + i\nu_\perp/b) \qquad \right\} = -i\nu_\parallel \frac{[\omega + i\nu_\parallel(\bar{\chi} - \bar{\xi})]}{[\omega + i\nu_\parallel\bar{\chi}]}$$

$$\cdot \left\{\omega[1 + \rho b] - \omega_D + i\nu_\perp\rho\right\} \qquad (1a)$$

neglecting gravity which has a minor effect on the drift mode and acoustic terms which are stabilizing for $k_z C_s \geq \omega_D$ (Hudson 1974). Here $L_\perp$ is the density gradient scale length, $b = k_y{}^2 C_s{}^2/\Omega^2$ is the ion FLR parameter, $\omega_D = -\sqrt{b}\, C_s/L_\perp$ is the diamagnetic drift frequency in terms of the ion acoustic speed $C_s{}^2 = T_e/M$, $\nu_\perp/b = \nu_{in} + 0.3\, b\, \nu_{ii}$ is the perpendicular ion diffusion rate, $\nu_\parallel = k_z{}^2 a_e{}^2/2C_r \nu_e$ is the parallel electron streaming rate, $\nu_e = \nu_{ei} + \nu_{es}$ is the total electron collision frequency, $\nu_{ji}$ and $\nu_{jn}$ ($j = 3, i$) are Coulomb and neutral collision frequencies $\rho = 1 + R_T = 1 + T_i/T_e$. We also define

$$\bar{\chi} = 2/3\left[C_r C_x + \left(1 + C_t\right)^2\right] + i\,2(m/M)\nu_e/\nu_\parallel \quad \cdot (m/M)\left(2\nu_{ei} + \nu_{en}\right)/\nu_\parallel$$

$$\bar{\xi} = 2/3\left(1 + C_t\right)^2 + 2\left(1 + C_t\right)\left(1 - R_t\right) \quad + i\ .\left(1 - R_T\right)(m/M)\left(\nu_{en} - 3\nu_{ei}\right)/\nu_\parallel$$

which are combinations of finite electron heat conduction terms, plotted in Figure 1 neglecting the energy transfer terms proportional to m/M which dominate at high densities and long parallel wavelengths ($k_z\lambda_e < \sqrt{m/M}$).

In the electrostatic dispersion relation (1a) the drift mode is coupled to the entropy mode (Tsai et al. 1970; Hudson and Kennel 1974b), interchange and acoustic terms. Finite parallel electron heat conduction is included in (1a), which reduces to the electrostatic dispersion relation obtained by

Tsai et al. (1970), when neutral collisions and energy transfer from electrons to ions and neutrals are neglected.

Isothermal Limit

The isothermal limit which applies to $k_z\lambda_e > 1$ is formally recovered from (1a) by letting the dimensionless coefficient of electron thermal conductivity $C_x \to \infty$. Hence $\bar{\chi} \to \infty$ in contrast to $\bar{\xi}$ which remains finite, and (1a) yields

$$b\left\{\left(\omega + R_T\omega_D\right)(\omega + i\nu_\perp/b) + g/L_\perp\right\}$$

$$= -i\nu_\parallel\{\omega[1 + \rho b] - \omega_D + i\nu_\perp\rho\} \tag{1b}$$

The collisionless limit $k_z\lambda_e > 1$ is only fully treated by a kinetic theory (Coroniti and Kennel 1970) which includes Landau resonance effects. Neglecting these resonances (1b) reduces to the fluid limit of their kinetic dispersion relation as $\nu_e \to 0$ or $\nu_\parallel \to \infty$.

In the opposite limit $\nu_\parallel \to 0$ for $R_T \sim 1$, we again find that $\bar{\chi} \to \infty$ while $\bar{\xi}$ remains finite. Hence by including energy transfer effects in $\bar{\chi}$ and $\bar{\xi}$, we have obtained the short parallel wavelength isothermal results at long parallel wavelengths. Including ion temperature fluctuations neglected in the derivation of (1a), we have found that this result is valid in a partially ionized plasma with $\nu_{in}/\nu_{ii} > 2\sqrt{m/M}$ at $T_e = T_i = T_n$ (Hudson 1974). For this small ratio of ion-neutral to Coulomb collision frequencies, the ion energy loss rate to the neutral sink exceeds the ion energy transfer

rate back to the electrons, and the plasma behaves isothermally at parallel wavelengths exceeding the electron energy transfer mean free path $\sqrt{M/m}\,\lambda_e$. However, in a fully ionized plasma there is a balance of energy exchange between ions and electrons, and the plasma behaves adiabatically at long parallel wavelengths.

The isothermal dispersion relation (1a) has therefore acquired a new significance at long parallel wavelengths where Landau resonance effects are negligible. Note that the cubic dispersion relation (1a) reduces to a quadratic (1b) with the elimination of the entropy mode in the isothermal approximation. It is useful first to obtain the marginal stability criterion in this isothermal approximation for the simple quadratic, since the same proceedure will be used to derive the general marginal stability criterion for the cubic.

## Isothermal Marginal Stability Criterion

At marginal stability $\omega$ is real and the real and imaginary parts of (1b) vanish separately. The imaginary part yields the dispersion relation

$$\omega = \omega_D \frac{\left(1 - R_T \nu_\perp/\nu_\parallel\right)}{\left(1 + \rho b + \nu_\perp/\nu_\parallel\right)}, \qquad \rho = 1 + R_T \tag{2a}$$

For $R_T \nu_\perp/\nu_\parallel << 1$ we obtain the drift mode dispersion relation

$$\omega = \omega_D/(1 + \rho\, b) \tag{2b}$$

The drift mode marginal stability criterion is obtained by substituting (2a) back into the real part of (1b), neglecting gravity

$$0 < \hat{K} = \frac{\left(1 - R_T \nu_\perp / \nu_\parallel\right)}{\left(1 + \rho b + \nu_\perp / \nu_\parallel\right)^2} \;.$$

$$K \equiv \nu_\perp \nu_\parallel \Big/ b \omega_D^2 \propto k_z^2 L_\perp^2 \tag{3}$$

Instability results for $0 < K < \hat{K}$, where K is the destabilizing parameter proportional to the density gradient scale length $L_\perp$. An unstable drift mode is only obtained for $R_T \nu_\perp / \nu_\parallel < 1$, which requires finite $k_z$. Equation (3) is plotted in Figure 2 for $R_T = 1$ and $b << \nu_\perp / \nu_\parallel < 1$. The threshold increases ($\hat{K} \propto k_z^2 L_\perp^2$ decreases) with increasing $\nu_\perp / \nu_\parallel$ proportional to total density.

General Marginal Stability Criterion

Having analyzed the marginally stable quadratic (1b) containing the FLR coupled interchange and drift modes in the isothermal limit, we now consider the marginally stable cubic (1a) containing the FLR coupled interchange, entropy and drift modes at intermediate parallel wavelengths $\sqrt{m/M} < k_z \lambda_e < 1$ where finite heat conduction effects apply in general.

The real and imaginary parts of (1a) vanish separately at marginal stability where $\omega$ is real. The imaginary part, which was linear in the isothermal case, is now quadratic with the addition of the entropy mode.

Expanding that quadratic assuming $|1 - R_T \nu_\perp/\nu_\parallel| > b\,(3R_T + K)$ yields the drift mode root

$$\frac{\omega_+}{\omega_D} = \frac{\left(1 - R_T\nu_\perp/\nu_\parallel\right) - R_T\bar{\chi}b}{\left[1 + \nu_\perp/\nu_\parallel + \left(1 + R_T + \bar{\chi}\right)b\right]} + \frac{(\chi - \xi)\left(1 + R_T\right)bK}{\left(1 - R_T\nu_\perp/\nu_\parallel\right)} \tag{4}$$

The instability criterion for the drift mode is obtained by substituting (4) back into the real part of (1a), yielding (Tsai et al. 1970)

$$\begin{aligned} K^{-1} - (1 + \nu_\perp/\nu_\parallel)^2 \frac{\left(1 + \bar{\xi} - R_T\nu_\perp/\nu_\parallel\right)}{\left(1 + R_T\nu_\perp/\nu_\parallel\right)^2} &> \\ b\Big\{(\bar{\chi} - \bar{\xi})^2 K\left[\nu_\parallel/\nu_\perp + 4R \quad 1 + \bar{\chi}\right] \\ - \bar{\chi}(\nu_\parallel/\nu_\perp)(\bar{\chi} - \bar{\xi}) + 4\left[4(\bar{\xi} - 1)\left(1 + R_T\right) + 1\right]\Big\} \\ + 0\left[b\nu_\perp/\nu_\parallel,\ b^2\right] \end{aligned} \tag{5}$$

This is the most general drift mode marginal stability criterion which applies to all densities $k_z\lambda_e < 1$.

We have plotted the marginally stable $\hat{K}$ vs. $\nu_\perp/\nu_\parallel$ in Figure 3 for a fully ionized plasma at $R_T = 1$ and several values of the parameter $N = 0.3\, L_\perp/\lambda_e$ proportional to density. We use the fully ionized values of $\bar{\chi} = 3.03$

and $\bar{\xi} = 1.95$ (Tsai et al. 1970) and substitute for $b = (1/N\sqrt{K\nu_\perp/\nu_\parallel}$. Comparison with the isothermal curve in Figure 2 demonstrates that the drift instability threshold is lower at low densities ($\nu_\perp/\nu_\parallel << 0.1$) than in the isothermal limit.

Analytic approximations to the high and low density limits can be obtained from the left and right hand sides of (5) respectively. At high densities the left hand side of (5) yields (Tsai et al. 1970)

$$K < \frac{\left(1 - R_T\nu_\perp/\nu_\parallel\right)^2}{\left(1 + \nu_\perp/\nu_\parallel\right)^2\left(1 + \bar{\xi} - R_T\nu_\perp/\nu_\parallel\right)} \tag{6}$$

which is plotted as the horizontal dashed curve in Figure 3. At low densities the right hand side of (5) yields

$$K < \frac{\bar{\chi}(\bar{\chi} - \bar{\xi})\nu_\parallel/\nu_\perp - 4\left[4(\bar{\xi} - 1)\left(1 + R_T\right) + 1\right]}{(\bar{\chi} - \bar{\xi})^2\left[\nu_\parallel/\nu_\perp + 4R_T + 1 + \bar{\chi}\right]} \tag{8a}$$

which approaches

$$K < \bar{\chi}/(\bar{\chi} - \bar{\xi}) \tag{8b}$$

as $\nu_\perp/\nu_\parallel \to 0$. Equation (7) is plotted as the vertical dashed curve in Figure 3. These approximate marginal stability criteria again demonstrate the increase at high densities and decrease at low densities, as compared to the isothermal threshold (3).

It is significant to note in Figure 3 that for all N values the exact marginal stability curves (5) cross at $\nu_\perp/\nu_\parallel = 0.1$, $K = 0.25$ where the approximate curves (6) and (7) cross. This single parameter point separates the finite heat conduction behavior of the drift mode into high and low density regions where the threshold is increased or decreased from its isothermal value.

In both the large and small $\nu_\perp/\nu_\parallel$ limits the finite heat conduction marginal stability criteria (6) and (7) reduce to the isothermal threshold (3) as $\bar{\chi} \to \infty$. In (6) $(1 + \bar{\xi})$ comes from $1 + \bar{\chi} - (\bar{\chi} - \bar{\xi})$, and $\bar{\xi}$ remains finite as we take the isothermal limit $\bar{\chi} \to \infty$. Hence for $k_z \lambda_e < \sqrt{m/M}$ the plasma behaves isothermally, independent of $\nu_\perp/\nu_\parallel$,, under the assumptions discussed in the preceeding section.

Drift Mode Growth Rate

Assuming $\omega/\nu_\parallel$, $\nu_\perp/\nu_\parallel$ and $bR_T\omega_D/\nu_\parallel << 1$ decouples the drift mode from entropy and interchange mode terms in (1a) and yields an approximate expression for the drift mode dispersion relation and growth rate in the low density finite heat conduction limit

$$\omega = \omega_D/(1 + \rho b) \tag{15a}$$

$$\gamma/\nu_\perp = \rho/(1 + \rho b)\left[\bar{\chi}/(\bar{\chi} - \bar{\xi})\left(K^{-1}/(1 + \rho b)\right) - 1\right] \tag{15b}$$

Neglecting FLR corrections, this result agrees with Ellis and Motley (1974). In the limit $\overline{\chi} \to \infty$ the isothermal growth rate is recovered (Chu et al. 1969). It is apparent that finite heat conduction increases the drift mode growth rate by the factor $\chi/(\chi - \xi) = (5/3)^2$ in a fully ionized plasma at low densities ($\nu_\perp/\nu_\parallel < 0.1$). At higher densities($\nu_\perp/\nu_\parallel > 0.1$) and well above marginal stability where the preceeding assumptions break down the drift mode growth rate must be obtained by solving the cubic dispersion relation (1a) in general.

The drift mode propagates in the electron diamagnetic drift direction. It is destabilized by the difference between the electron guiding center current and the ion current which is an ion FLR effect of order b. In contrast, in the companion paper we have seen that the interchange mode is stabilized by FLR effects in a weakly collisional plasma($\nu_\perp/b < 2\sqrt{g/L_\perp}$). We therefore expect the drift mode growth rate to peak at shorter perpendicular wavelengths (Hudson and Kennel 1974c). Damping is proportional to the perpendicular ion diffusion rate $\nu_\perp/b = \nu_{in} + 0.3\, b\nu_{ii}$.

## Conclusions

We have obtained a marginal stability criterion for the drift mode in a partially ionized plasma which is valid for all wavelengths greater than the total electron mean free path ($k_z\lambda_e < 1$). An apparent contradiction in the effect of finite heat conduction on the instability threshold and growth rate has been resolved (Tsai et al. 1970; Ellis and Motley, 1974). At low densities ($\nu_\perp/\nu_\parallel < 0.1$) the drift mode decouples from the entropy mode

and the growth rate of the former is enhanced. At higher densities ($\nu_\perp/\nu_\parallel > 0.1$) the drift and entropy modes are coupled, and the drift mode growth rate is suppressed.

When $k_z\lambda_e < \sqrt{m/M}$ the drift mode approaches the isothermal electron limit for all $\nu_\perp/\nu_\parallel$, and we recover the isothermal instability criterion of Chu et al. (1969) generalized to a partially ionized plasma. This assumes that ions behave isothermally due to energy loss to the neutral sink, which Hudson (1974) has shown is valid for $\nu_{in}/\nu_{ii} > 2\sqrt{m/M}$ at $T_e = T_i = T_n$. In a fully ionized plasma there is a balance of energy exchange between ions and electrons in long parallel wavelength modes ($k_z\lambda_e < \sqrt{m/M}$) which behave adiabatically. This distinction in behavior at long parallel wavelengths is significant since the drift instability threshold and growth rate is modified by the nonisothermal effects considered above.

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Figure Captions

Figure 1. Plot of energy heat flow parameters for $R_T = 1$ neglecting energy transfer terms. $\chi = 2/3\,[C_r C_x + (1+C_t)^2]$ and $\xi = 2/3\,(1+C_t)^2$, where $C_r$, $C_t$, and $C_x$ are the dimensionless resistive, thermal force, and thermal conductivity coefficients computed by Braginskii (1965) for a fully ionized plasma: $C_r = 0.15$, $C_t = 0.71$ and $C_x = 3.16$. Schunk and Walker (1970) have plotted them for a partially ionized plasma as $g_\sigma = C_r$, $g_\sigma/g_\mu = C_t^{-1}$ and $(1/g_k - 0.5\, g_\sigma/g_\mu g_r) = C_x/5$ from Shkarofsky's (1961) tabulation of the g-factors. Here we assume that electrons collide with hard sphere neutrals at a rate proportional to the electron thermal speed $v = a_e$.

Figure 2. Plot of the isothermal drift mode marginally stable parameter $K \equiv \nu_\perp \nu_\parallel / b\omega_D^2$ as a function of $\nu_\perp/\nu_\parallel$ from (3) at $R_T = 1$ for $b << \nu_\perp/\nu_\parallel < 1$. The threshold increases ($\hat{K} \propto k_z^2 L_\perp^2$ decreases) with increasing $\nu_\perp/\nu_\parallel$ proportional to total density.

Figure 3. Plot of the fully ionized finite heat conduction drift mode marginally stable parameter $K \equiv \nu_\perp \nu_\parallel / b\omega_D^2$ as a function of $\nu_\perp/\nu_\parallel$ from (5) at $R_T = 1$ for several values of $N = 0.3\, L_\perp/\lambda_e$ proportional to density. The exact curves are well approximated by the dashed lines (6) and (7) at high and low densities; they correspond to the left and right hand sides of (5) which are separately independent of N. The crossover point at $\nu_\perp/\nu_\parallel = 0.1$ corresponds to the separation between the low and high density finite heat conduction regimes where the drift instability threshold is lowered ($\hat{K} \propto k_z^2 L_\perp^2$ increased) and raised compared to the isothermal threshold in Figure 2.

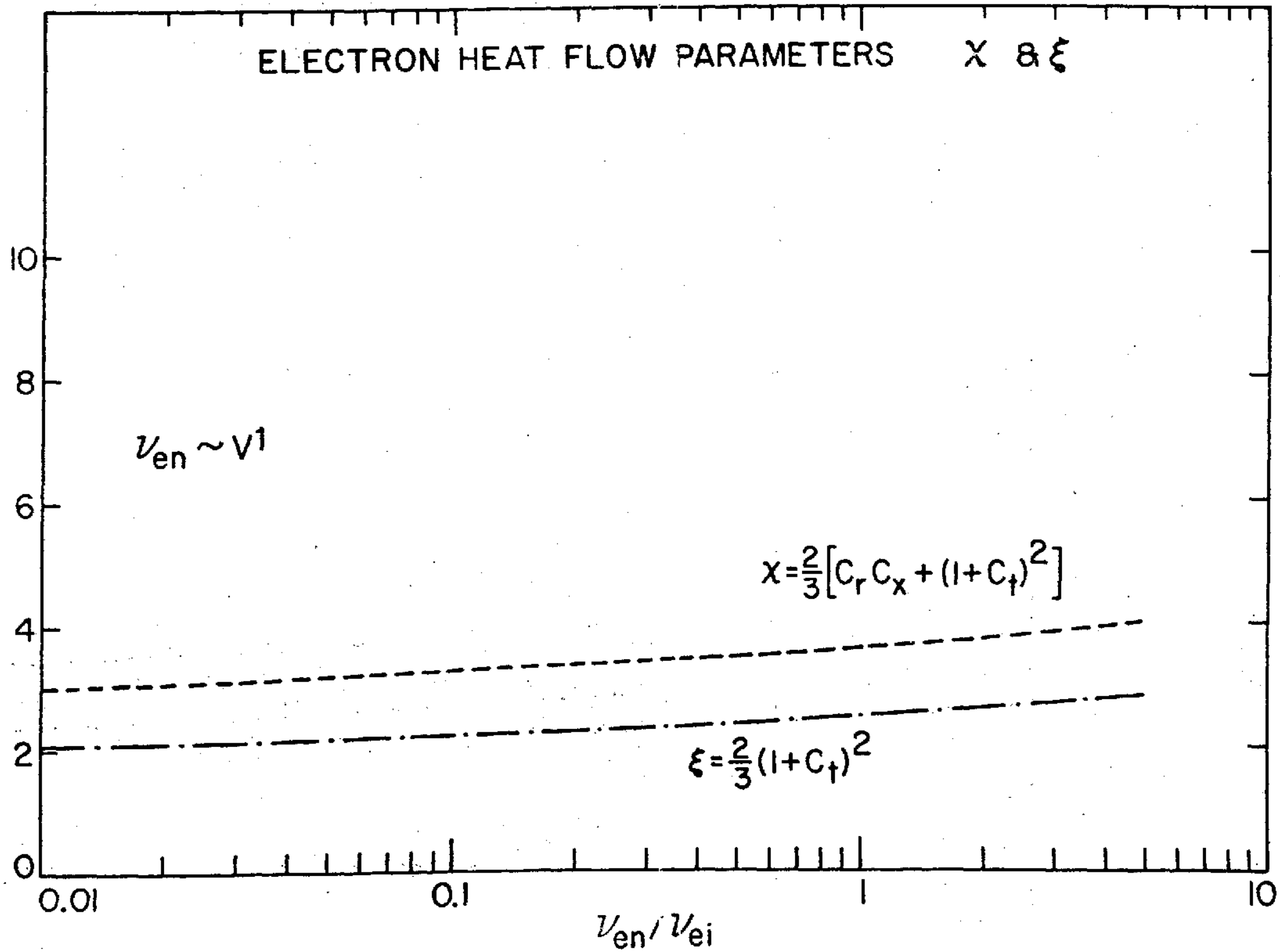

ELECTRON HEAT FLOW PARAMETERS $\chi$ & $\xi$
$\nu_{en} \sim V^1$
$\chi = \frac{2}{3}\left[C_r C_x + (1+C_t)^2\right]$
$\xi = \frac{2}{3}(1+C_t)^2$
0
2
4
6
8
10
0.01
0.1
1
10
$\nu_{en}/\nu_{ei}$

Figure 1

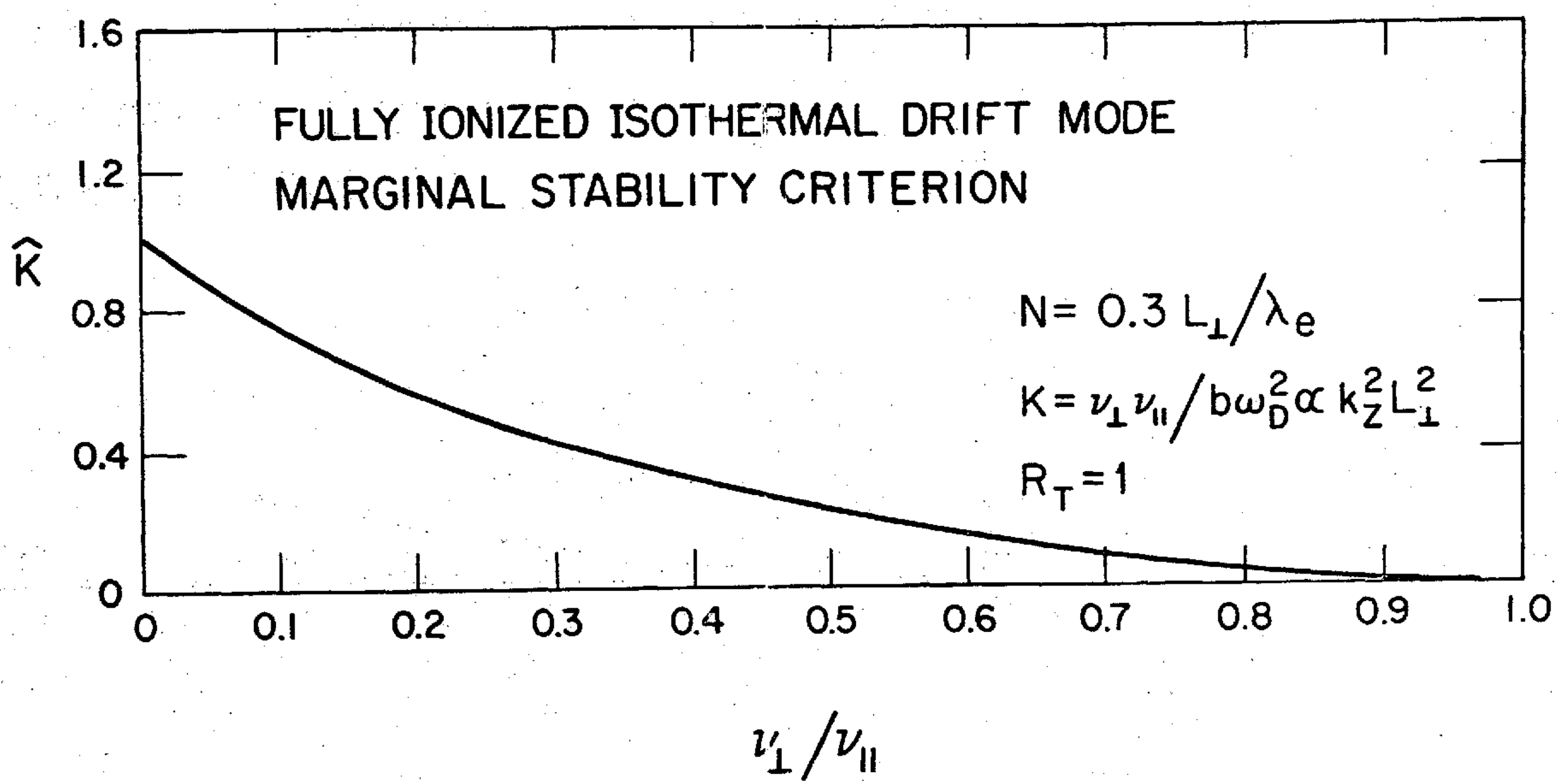
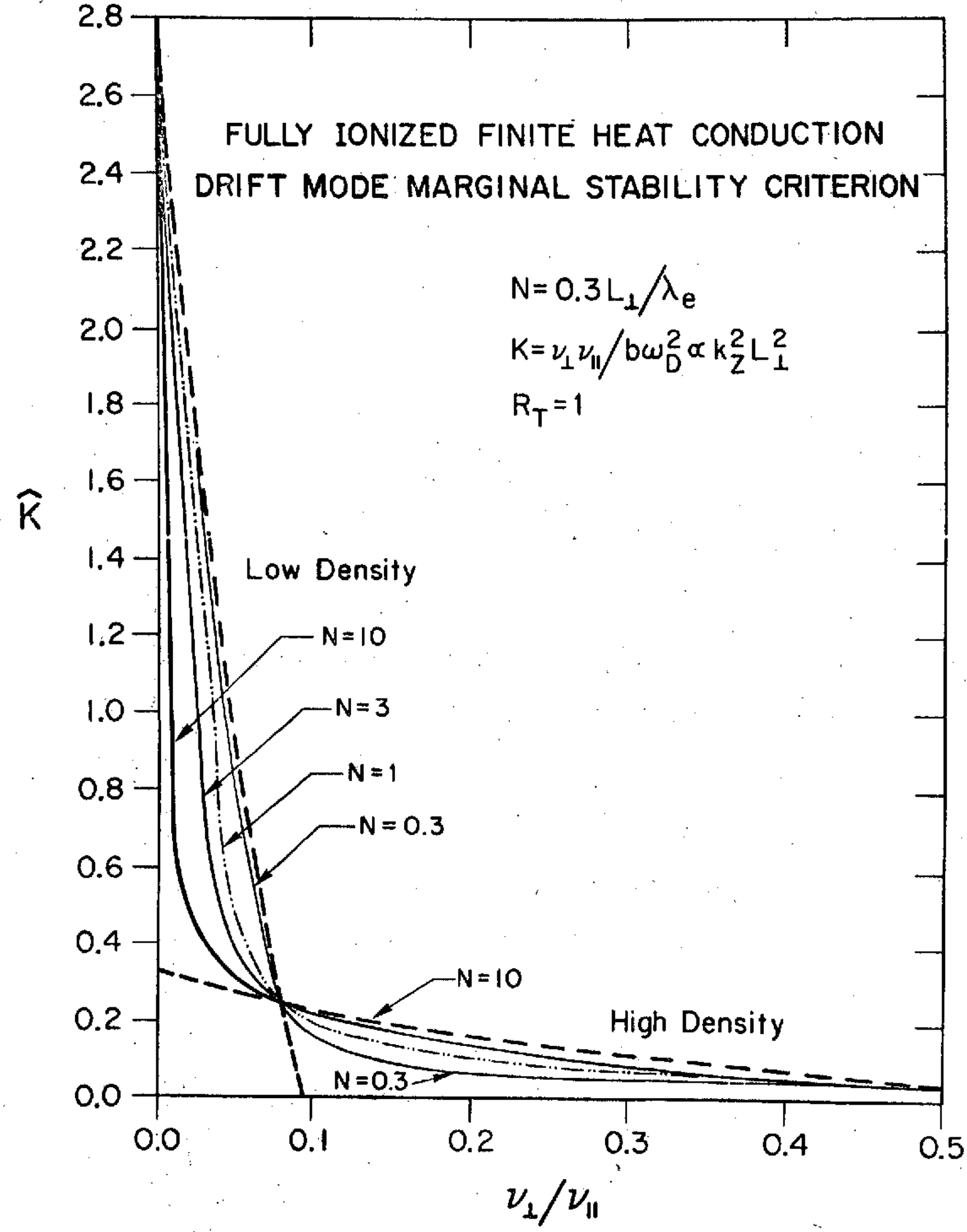

FULLY IONIZED ISOTHERMAL DRIFT MODE
MARGINAL STABILITY CRITERION
$N = 0.3\, L_\perp / \lambda_e$
$K = \nu_\perp \nu_\parallel / b\omega_D^2 \propto k_Z^2 L_\perp^2$
$R_T = 1$
$\hat{K}$
1.6
1.2
0.8
0.4
0
0
0.1
0.2
0.3
0.4
0.5
0.6
0.7
0.8
0.9
1.0
$\nu'_\perp / \nu_\parallel$

Figure 2



FULLY IONIZED FINITE HEAT CONDUCTION
DRIFT MODE MARGINAL STABILITY CRITERION
$N = 0.3\, L_\perp / \lambda_e$
$K = \nu_\perp \nu_\parallel / b\omega_D^2 \propto k_Z^2 L_\perp^2$
$R_T = 1$
Low Density
N=10
N=3
N=1
N=0.3
N=10
High Density
N=0.3
$\hat{K}$
$\nu_\perp / \nu_\parallel$
2.8
2.6
2.4
2.2
2.0
1.8
1.6
1.4
1.2
1.0
0.8
0.6
0.4
0.2
0.0
0.0
0.1
0.2
0.3
0.4
0.5


Figure 3

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